Lepton Flavor Violation and Supersymmetric Dirac Leptogenesis

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Abstract

Dirac leptogenesis (or Dirac neutrinogenesis), in which neutrinos are purely Dirac particles, is an interesting alternative to the standard leptogenesis scenario. In its supersymmetric version, the modified form of the superpotential required for successful baryogenesis contributes new, generically non-flavor-diagonal terms to the slepton and sneutrino mass matrices. In this work, we examine how current experimental bounds on flavor-changing effects in the lepton sector (and particularly the bound on $\mu \to e\gamma$) constrain Dirac leptogenesis and we find that it is capable of succeeding with superpartner masses as low as ~ 100 GeV. For such light scalars and electroweakinos, upcoming experiments such as MEG are generically expected to observe signals of lepton flavor violation.

1 Introduction

When singlet fermions are not present in a given theory, Dirac leptogenesis [1, 2], or Dirac neutrinogenesis, represents a very interesting alternative to the traditional leptogenesis scenario, which relies on the existence of heavy Majorana neutrinos. In Dirac leptogenesis, neutrinos are purely Dirac particles whose small but nonzero masses appear as ratios of dimensionful parameters in an effective field theory. It has been shown [3] that, in the context of split supersymmetry [4, 5], Dirac leptogenesis is a phenomenologically viable scenario capable of satisfying all relevant constraints from cosmology and neutrino physics as well as reproducing the observed baryon-to-photon ratio η of the universe. Split supersymmetry is advantageous primarily for two reasons. The first of these is that very heavy gravitinos can easily evade the constraints that big bang nucleosynthesis (BBN) places on the post-inflationary reheating temperature; the second is that in Dirac leptogenesis the superpotential is extended by the addition of new terms carrying phases and nontrivial flavor structure. Dangerous contributions from these terms to flavor-changing processes are automatically safe in split supersymmetry, but must be treated with care when the scale of all the superpartners is near the electroweak scale. In the latter case (assuming that the constraints for gravitino cosmology are satisfied) it is necessary to compute carefully the rates for flavor-violating processes in the lepton sector.

The aim of this paper is to ascertain whether or not Dirac Leptogenesis is permitted by flavor violation constraints in supersymmetric theories with low-scale slepton masses. We will begin our inquiry by briefly reviewing the theoretical framework of Dirac leptogenesis and deriving the additional contributions to the slepton mass matrices to which its superpotential gives rise. We then turn to the calculation of rates for the processes $\mu \to e\gamma$ and $\tau \to \mu\gamma$. Finally, we apply the combined constraints from flavor violation and baryogenesis and assess the viability and predictability of Dirac leptogenesis when superpartner masses are at or around the weak scale.

2 Dirac Leptogenesis and the Slepton Mass Matrices

In Dirac leptogenesis, as in the traditional leptogenesis picture [6, 7], the conditions for successful baryogenesis [8] are met by positing the existence of a heavy particle with CP-violating decays into leptons.[1] The lepton number produced in these decays is then processed into a baryon number for the universe by electroweak sphaleron processes [11]. In the supersymmetric version of Dirac leptogenesis, all of this is engineered via a specific set of modifications to the superpotential and the postulation of a few additional superfields. In addition to the matter fields of the MSSM, at least two massive vector-like pairs of chiral superfields Φ and $\overline{\Phi}$ are required, as are three generations of right-handed neutrino superfield N_a and an additional exotic superfield χ, whose function will be to acquire a scalar VEV. An additional symmetry, whose breaking will be responsible for late neutrino masses, is also imposed, and charges under it assigned so that the most general leptonic-sector superpotential is

$$\mathcal{W} \ni \lambda_{i\alpha} N_\alpha \Phi_i H_u + h_{i\alpha} L_\alpha \overline{\Phi}_i \chi + M_{\Phi_i} \Phi_i \overline{\Phi}_i + \mu H_u H_d. \tag{1}$$

Here, λ and h are (generally complex) coupling matrices and M_{Φ_i} are the masses of the heavy vector-like pairs, which, in order for successful baryogenesis to occur, are generically required to be 10^{10} GeV or larger.[2] Once the heavy superfields Φ_i and $\overline{\Phi}_i$ are integrated out, the resulting effective superpotential

$$\mathcal{W}_{eff} \ni \frac{\lambda_{i\alpha} h^*_{i\beta}}{M_{\Phi_i}} \chi L_\beta H_u N_\alpha + \mu H_u H_d \tag{2}$$

will yield a small but nonzero Dirac mass term for the neutrinos, provided that some mechanism is invoked to give the scalar component of χ a VEV. Since λ and h are complex, the decays of both the scalar and fermionic components of Φ_1 and $\overline{\Phi}_1$, the lightest of the additional heavy superfields, will in general be CP-violating[3] and will result in two equal and opposite stores of lepton number, L_{agg} and L_R. The first of these is an aggregate lepton number stored in left-handed leptons, sleptons,

[1] For recent variations on both supersymmetric and non-supersymmetric scenarios, see for example [9, 10].

[2] In general M_Φ, λ, and h are all complex matrices, but we can always choose to work in a basis where M_Φ is diagonal.

[3] It can be shown that there is at least one nontrivial CP-violating phase in λ and h.

and other fields in equilibrium with them, and is transformed into baryon number by sphaleron interactions. The second is stored only in right-handed neutrinos,[4] which, being singlets under $SU(2) \times U(1)_Y$, do not experience sphaleron effects and only couple to the other light matter fields through the effective neutrino Yukawa interaction given by (2) with $\chi \to \langle \chi \rangle$. This interaction is suppressed by $\langle \chi \rangle / M_{\Phi_1}$ and therefore the time scale for the equilibration of left-handed and right-handed stores of lepton number can be quite late. If the effective neutrino Yukawas are sufficiently small,[5] the equilibration time scale will be much longer than the time scale at which sphaleron processes effectively shut off [1]. When this is the case, a net baryon number for the universe will already have frozen in and will persist unaltered until present time.

In addition to providing a mechanism for baryogenesis, Dirac leptogenesis holds some interesting implications for neutrino physics. Here, the squared neutrino masses are given by

$$|m_\nu|^2_{\alpha\beta} = \Big(v \langle \chi \rangle \sin\beta \Big)^2 \sum_{i,j=1}^{2\ (or\ 3)} \sum_{\gamma=1}^{3} \lambda^*_{i\gamma} \lambda_{j\gamma} h^*_{i\alpha} h_{j\beta} \frac{1}{M^*_{\Phi_i} M_{\Phi_j}}. \tag{3}$$

In order to reproduce the observed mass splittings [12]

$$\Delta m^2_{21} = \left(7.9^{+0.6}_{-0.6}\right) \times 10^{-5} \mathrm{eV}^2, \qquad |\Delta m^2_{31}| = \left(2.2^{+0.7}_{-0.5}\right) \times 10^{-3} \mathrm{eV}^2, \tag{4}$$

where $\Delta m^2_{ij} \equiv m^2_i - m^2_j$, and the angles in the U_{MNS} matrix, it is necessary to impose a few conditions on the coupling matrices λ and h, which determine the matrix structure in (3). There are several ways of doing this, but we will focus on one particular, theoretically-motivated scenario called constrained hierarchical Dirac leptogenesis (CHDL) [3], which produces a normal hierarchy among neutrino masses.

In CHDL, the appropriate matrix structure for λ, h, and the (diagonal) mass matrix M_Φ is obtained by requiring that λ and h both be antisymmetric with $\mathcal{O}(1)$ entries and by relating M_Φ to the fermion Yukawas through a flavor symmetry. This structure ensures that the neutrino mass matrix one obtains after integrating out the heavy fields Φ_i and $\overline{\Phi}_i$ corresponds to a normal neutrino hierarchy, provided that there is a hierarchy among the mass eigenstates M_{Φ_1} and M_{Φ_2}. Of course this coupling structure also impacts baryogenesis, but it has been shown that in CHDL, the constraints from baryogenesis and neutrino physics can be satisfied simultaneously [3].

Baryogenesis and a small neutrino Dirac mass are not the only consequences of equation (1), however. Assuming that at some high scale M all the soft supersymmetry breaking terms are flavor diagonal and universal, the new terms of equation (1) will induce potentially off-diagonal contributions to the slepton mass matrices m^2_{LL}, m^2_{RR} and m^2_{LR}. As in Majorana leptogenesis some of these contributions come from running the scalar soft masses from some high scale M (the Planck scale, the GUT scale, etc.) down to M_{Φ_1}. Assuming the Universality condition at the high scale and that all soft A-terms are equal to the relevant yukawa coupling multiplied by the

[4] We assume that right-handed sneutrinos couple strongly enough to other fields in the theory (e.g. with left-handed sneutrinos in the event that χ has a large F-term VEV) to equilibrate with them and thus contribute to L_{agg}. Otherwise they would contribute to L_R.

[5] Recall that in this scenario small (Dirac) neutrino masses are obtained through small (effective) neutrino Yukawas.

universal soft supersymmetric mass m_s, one can simply estimate the flavor violating corrections to the mass matrix by integrating the RGE equations iteratively [13, 14], and generally obtain off-diagonal contributions to the slepton masses δm^2_{LL} and δm^2_{LL}:

$$\begin{aligned} \delta m^2_{LL} &\approx -\frac{1}{2\pi^2}\ln\left(\frac{M}{M_{\Phi_1}}\right) h^*_{i\alpha}h_{i\beta}m^2_s && (5) \\ \delta m^2_{RR} &\approx -\frac{1}{2\pi^2}\ln\left(\frac{M}{M_{\Phi_1}}\right) \lambda^*_{i\alpha}\lambda_{i\beta}m^2_s. && (6) \end{aligned}$$

If the F-term of χ acquires a VEV $\langle F_\chi\rangle$ (a corollary in most mechanisms via which its scalar component obtains a VEV), the effective theory superpotential yields yet another potentially off diagonal scalar mass term

$$\delta m^2_{LR} = h^\dagger_{i\alpha}\lambda_{i\beta}\frac{\langle F_\chi\rangle}{M_{\Phi_1}} v\sin\beta \tag{7}$$

after electroweak symmetry breaking, which mixes left-handed and right-handed sneutrinos.

In CHDL, where we have a specific flavor structure for the matrices λ, h and M_Φ, one has a specific prediction for flavor mixing among sleptons once the electroweak and hidden symmetries are broken. In order to examine the effect of these mixings, the full mass matrices for both the charged sleptons and sneutrinos must be taken into account. For simplicity, we will continue to assume that the leading soft breaking sector is flavor diagonal and universal with a common scalar mass m_s. The resulting additional contributions to the slepton mass squared matrices, given by equations (5), (6), and (7), can thus be expressed in terms of the 3×3 submatrices δm^2_{LL}, δm^2_{RR}, and δm^2_{LR} as

$$\delta m^2_{\tilde\ell^\pm} = \left(\begin{array}{c:c} \delta m^2_{LL} & 0 \\ \hdashline 0 & 0 \end{array}\right) \qquad \delta m^2_{\tilde\nu} = \left(\begin{array}{c:c} \delta m^2_{LL} & \delta m^2_{LR} \\ \hdashline (\delta m^2_{LR})^\dagger & \delta m^2_{LL} \end{array}\right). \tag{8}$$

The only contribution to the charged slepton mass squared matrix comes from δm^2_{LL}, while the sneutrino mass squared matrix receives not only additional flavor mixings among left-handed and among right-handed sneutrinos, but also an effective A-term from δm^2_{LR} which intermixes left-handed and right-handed sneutrinos.

Once the matrices λ and h have been fixed, up to an overall scaling parameter[6] f, to give the correct leptonic mixing matrix U_{MNS}, the remaining free parameters of the model are $\langle\chi\rangle$, M_{Φ_i}, and f. After the mass scale of the neutrinos is fixed, these parameters are not independent anymore, and are related by the constraint embodied in equation (3). In CHDL, where the lightest

[6] In CHDL λ and h are antisymmetric, with $\mathcal{O}(1)$ entries up to an overall scaling parameter f

$$\lambda \approx f\begin{pmatrix} 0 & 1 & 1 \\ -1 & 0 & a_3 \\ -1 & -a_3 & 0 \end{pmatrix} \qquad h \approx f\begin{pmatrix} 0 & 1 & 1 \\ -1 & 0 & b_3 \\ -1 & -b_3 & 0 \end{pmatrix}, \tag{9}$$

where the values of a_3 and b_3 depend on the hierarchy between M_{Φ_1} and M_{Φ_2}: when $M_{\Phi_2}/M_{\Phi_1} = m_\mu/m_e$, b_3 can vary between 1.4 and 2.9, while a_3 can vary between 35 and 90; when $M_{\Phi_2}/M_{\Phi_1} = 10$, b_3 can vary between 1.4 and 2.9, while a_3 can vary between 1.5 and 4.5 [3]).

neutrino is very light relative to the other two, the masses of the two heavier neutrinos are

$$
\begin{aligned}
m_{\nu_2}^2 &\approx \Delta m_{21}^2 = \left(7.9^{+0.6}_{-0.6}\right) \times 10^{-5} \mathrm{eV}^2 && (10) \\
m_{\nu_3}^2 &\approx \Delta m_{31}^2 = \left(2.2^{+0.7}_{-0.5}\right) \times 10^{-3} \mathrm{eV}^2. && (11)
\end{aligned}
$$

Plugging in these values and using $a_3 = 4.5$ and $b_3 = 2.2$ (the values most advantageous for baryogenesis with a chosen hierarchy $M_{\Phi_2}/M_{\Phi_1} = 10$) in the CHDL parametrization (9) of λ and h, we find that the relation between the parameters $\langle\chi\rangle$, M_{Φ_i}, and f, obtained from the neutrino mass, is

$$
\frac{f^2 \langle\chi\rangle}{M_{\Phi_1}} \sin\beta = 1.009 \times 10^{-13}. \tag{12}
$$

Using this relation, we can express the overall dependence of the slepton mass terms in (5), (6), and (7) on the relevant mass scales in the theory.

$$
\begin{array}{cccccc}
\delta m_{LL}^2 & \propto & f^2 & \propto & \frac{M_{\Phi_1}}{\langle\chi\rangle} \\
\delta m_{RR}^2 & \propto & f^2 & \propto & \frac{M_{\Phi_1}}{\langle\chi\rangle} \\
\delta m_{LR}^2 & \propto & \frac{f^2}{M_{\Phi_1}} & \propto & \frac{1}{\langle\chi\rangle}.
\end{array} \tag{13}
$$

Note that the proportionality constants for the bottom two equations are not dimensionless: the ones associated with δm_{LL}^2 and δm_{RR}^2 each contain a factor of m_s^2 and have mass dimension $[m]^2$, while the one associated with δm_{LR}^2 contains a factor of $\langle F_\chi\rangle v$ and has mass dimension $[m]^3$.

3 Flavor Violation

The most stringent constraints on flavor violation in the lepton sector come from measurements of the branching ratios for flavor-violating decays and conversions of heavy leptons, such as $\mu \to e\gamma$, $\tau \to \mu\gamma$, $\mu \to eee$ and $\mu A \to eA$. The current experimental limits on the 2-body decay processes are [15]

$$
\begin{aligned}
BR(\mu \to e\gamma) &< 1.2 \times 10^{-11}, && (14) \\
BR(\tau \to \mu\gamma) &< 1.1 \times 10^{-6}. && (15)
\end{aligned}
$$

In the near future, the MEG experiment [16] is expected to improve the current experimental bound on $\mu \to e\gamma$ by several orders of magnitude, to $\mathcal{O}(10^{-13} - 10^{-14})$ or lower. Other related projects, such as PRIME [17] (sensitive to $\mu A \to eA$ conversion), are expected to go online over the next few years. Projects have also been proposed [18] that would lower the bound from $\tau \to \mu\gamma$ to $\mathcal{O}(10^{-9})$.

The effective interaction leading to lepton flavor violating decays of the form $\ell_i \to \ell_j\gamma$, where ℓ_i and ℓ_j are charged leptons, can be written as

$$
\mathcal{I} = i\, e\, m_{\ell_j}\, \bar{u}_i(q-p)\, \sigma_{\alpha\beta}\, q^\beta \left(A^L P_L + A^R P_R\right) u_j(p)\, \epsilon^*(q), \tag{16}
$$

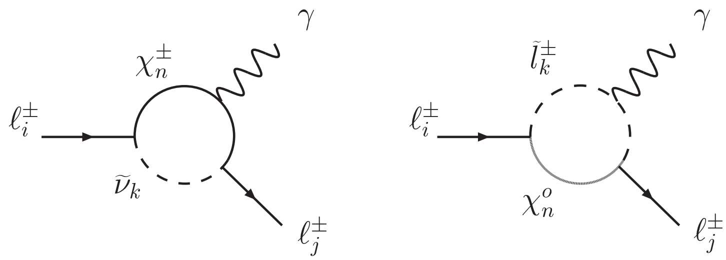

Figure 1: Feynman diagrams giving the two leading order contributions to the flavor-changing process $\ell_i^- \to \ell_j^- \gamma$ due to sneutrino (left diagram) and charged slepton (right diagram) mass mixings.

where q and p are the momenta of the photon and the outgoing lepton ℓ_j respectively, and m_{ℓ_j} is the outgoing lepton mass. The resulting decay rate is

$$\Gamma(l_j^- \to l_i^- \, \gamma) = \frac{e^2}{16\pi} m_{l_j}^5 (|A^L|^2 + |A^R|^2). \tag{17}$$

The leading contributions to the amplitudes A^L and A^R appear at one loop level and are shown in figure 1. They involve both a sneutrino (and chargino) mass eigenstate and charged slepton (and neutralino) mass eigenstate running in the loop. These amplitudes have been computed in [13] for a general MSSM scenario[7] but for completeness we include them in Appendix A.

In order to proceed further, it will be necessary to make a few assumptions concerning the supersymmetric model parameters. Those relevant to a discussion of lepton-sector flavor violation include the gaugino masses M_1 and M_2, the Higgs mass parameter μ, the ratio of Higgs VEVs $\tan\beta$, and the soft masses for the sleptons. In our analysis, we choose the values $M_1 = 160$ GeV, $M_2 = 220$ GeV, $\mu = 260$ GeV, and $\tan\beta = 3, 10$ and 30. As for the slepton soft masses we will, as previously mentioned, assume a common scale m_s and examine what effect varying m_s has on $BR(\mu \to e\gamma)$ and $BR(\tau \to \mu\gamma)$. We will assume that the scale at which soft masses are universal is $M = 2 \times 10^16$ GeV, though the results are not particularly sensitive to this choice.

The results of our calculation are displayed in figure 2. In the left panel, we show exclusion contours in M_{Φ_1}-$\langle\chi\rangle$ space for $m_s = 200$ GeV. The areas below and to the right of the lower contour (the white region) are excluded by the experimental bounds given in (14). We also include contours demarcating the region wherein baryogenesis can succeed, which have been updated from [3] to include the effects[8] of processes second order in Φ_1 and $\overline{\Phi}_1$ (see Appendix B). Contours have also been computed for $\tau \to \mu\gamma$, but the constraints they imply for the theory are far weaker than those from $\mu \to e\gamma$. In the right panel, we show how varying the universal scalar mass affects the branching ratio for $\mu \to e\gamma$, which reaches a maximum when m_s is around the weak scale. This is to be expected: when m_s is much larger than the weak scale both the slepton mass-squared

[7] In our case we need to add three right handed sneutrinos, but it is trivial to extend the result to include six sneutrino mass eigenstates instead of three.

[8] We thank A. Strumia for pointing out to us the potential numerical importance of these thermalization processes involving gauge interactions[19, 20].

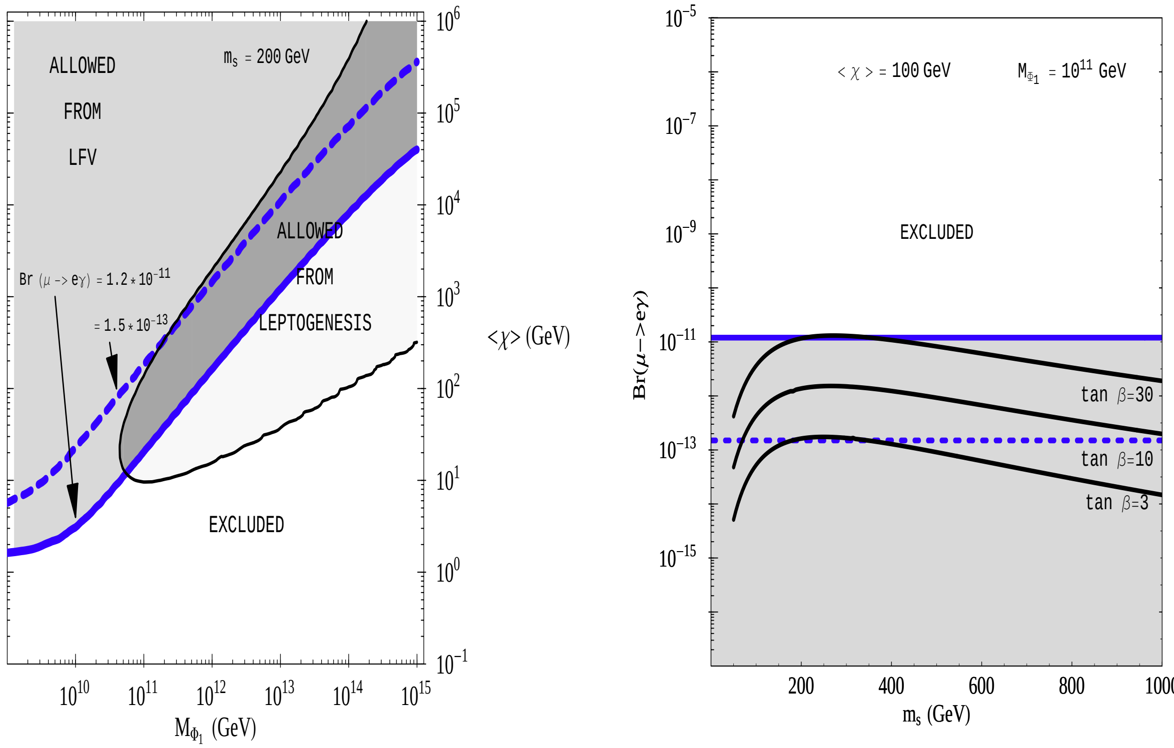

Figure 2: Exclusion plots combining constraints from both leptogenesis and flavor violation in the process $\mu \to e\gamma$. The left-hand panel shows exclusion contours in M_{Φ_1}-$\langle\chi\rangle$ space for a universal scalar soft mass $m_s = 200$ GeV, with $\tan\beta = 10$; the right hand panel shows the variation of the branching ratio $BR(\mu \to e\gamma)$ with respect to m_s using $\tan\beta = 3, 10$ and 30. In both plots, we have taken $M_1 = 160$ GeV, $M_2 = 220$ GeV, and $\mu = 260$ GeV. We have also assigned the χ superfield an F-term VEV $\sqrt{\langle F_\chi\rangle} = 10^7$ GeV. Such a large VEV results in large trilinear couplings between Higgs fields and sneutrinos and therefore induces potentially sizeable mixings between left-handed and right-handed sneutrinos after electroweak symmetry breaking. In each plot, the thick solid contours represent the current experimental bound on the branching fraction (14); the dashed lines represent the expected future experimental bound of 1.5×10^{-13} from MEG. The thin solid contour in the left-hand panel delimits the region allowed by leptogenesis constraints.

eigenvalues and the flavor-violating terms scale like m_s^2 and the sneutrino and charged slepton mixing matrices asymptote to a constant value, while the branching ratio is still suppressed by the masses running in the loop; as m_s decreases below the weak scale, δm^2_{LL} and δm^2_{LL} go to zero and the slepton masses are dominated by flavor diagonal electroweak contributions. We also observe that, as in the SUSY see-saw case [13, 14], the flavor violation rate is quite sensitive to $\tan\beta$.

In interpreting the results in figure 2, it is useful to note that in the regions of the plot near the exclusion contours (where the branching ratio for $\mu \to e\gamma$ is quite low), flavor-violating effects will be small. It is therefore valid to use the mass-insertion approximation there and treat δm^2_{LL} (left), δm^2_{LR}, and δm^2_{RR} as small corrections to the slepton propagators. Since the sneutrino propagator can receive mass insertions from all three, we will focus our analysis on the partial amplitude from the sneutrino-mediated process (the left diagram in figure 1). In figure 3, we list the leading contributions to this partial amplitude involving each of δm^2_{LL} (left diagram), δm^2_{LR}, and δm^2_{RR}. Since there is no coupling between leptons and right-handed sneutrinos, corrections from δm^2_{LR} and δm^2_{RR} only appear at second and third order in the mass insertion expansion, respectively. Therefore, if there is no substantial hierarchy among these three sets of mixing terms, the primary source of flavor-violation comes from mixings between left-handed sleptons. In the mass insertion approximation, the branching ratio for such processes can naively be estimated as

$$BR(\mu \to e\gamma) \propto \frac{\alpha^3}{G_F^2}\frac{(\delta m^2_{LL})^2}{m_s^8} \tag{18}$$

and therefore contours of branching ratio in the $M_{\Phi_1} - \langle\chi\rangle$ plane correspond to contours of δm^2_{LL}. According to equation (13), $\delta m^2_{LL} = c_1 M_{\Phi_1}/\langle\chi\rangle$, where c_1 is a dimensionless proportionality constant with dimension $[m]^2$, so the exclusion contour associated with left-left mixing takes the form

$$\ln M_\Phi = \ln\langle\chi\rangle + C_{LL}, \tag{19}$$

where $C_{LL} = -\ln(\delta m^2_{LL}/c_1)$. The oblique, upper exclusion contour in figure 2, which embodies the constraint from left-left mixing, is associated with this linear equation.

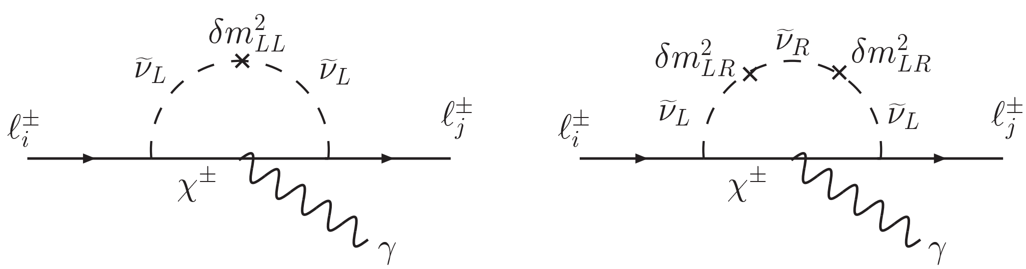

Figure 3: Feynman diagrams for the leading-order process involving δm^2_{LL} (left diagram), and for the leading process involving δm^2_{LR} (right diagram), with sneutrinos running in the loop in the mass-insertion approximation. Note that any process involving δm^2_{LR} necessarily involves two mass insertions, and any one involving δm^2_{RR} (given by the diagram on the right with an additional δm^2_{RR} insertion) necessarily involves three.

While some hierarchy among the mass insertion terms is necessary for δm^2_{LR} and δm^2_{RR} to be relevant, there is no a priori reason why such a hierarchy should not exist. The δm^2_{LR} contribution (7) is proportional to $\langle F_\chi\rangle$, which is essentially a free parameter. As mentioned above, $\langle F_\chi\rangle$ is not even relevant to Dirac leptogenesis per se, but appears as a common side-effect of mechanisms for assigning the χ superfield a scalar VEV. Still, in many such mechanisms [21], the scale $\sqrt{\langle F\rangle}$ can potentially be quite large (10^6 GeV or higher), and if this is the case, contribution from δm^2_{LR} could potentially be as important as those from δm^2_{LL}. Let us assume for a moment that this is the case and examine the constraints related to δm^2_{LR} and δm^2_{LR} together. In regions of figure 2 where $\langle\chi\rangle$ is small, we now have

$$BR(\mu \to e\gamma) \propto \frac{\alpha^3}{G_F^2 m_S^8} \frac{(\delta m_{LR}^2)^4}{m_{\tilde{\nu}_R}^4} = \text{constant} \tag{20}$$

along any exclusion contour. Equation (13) tells us that,

$$\delta m_{LR}^2 = c_2 \frac{1}{\chi}, \tag{21}$$

where c_2 has mass dimension $[m]^3$. Therefore, when $\delta m_{RR}^2 \ll m_{\tilde{\nu}_R}^2$, the associated contour is given by

$$\ln\langle\chi\rangle = C_{LR}, \tag{22}$$

where $C_{LR} = \ln(c_2/\delta m_{LR}^2)$. This equation explains the behavior of the contour in the left panel of figure 2 when $\langle\chi\rangle$ is small and the oblique bound from left-left mixing abruptly ells into a horizontal line—the bound from left-right mixing. In general, the effect of increasing $\langle F_\chi\rangle$ is to push this latter bound upward, and for $\sqrt{\langle F_\chi\rangle} \gtrsim 10^9$, the entirety of parameter space is excluded. The effect of right-right mixing is higher order still and only becomes relevant in regions of parameter space where M_{Φ_1} is large and $\langle\chi\rangle$ is small—regions already excluded by left-left mixing.

For the numerical analysis shown in Figure 2, we have taken $\sqrt{\langle F_\chi\rangle} = 10^7$ GeV which in some parts of the $\langle\chi\rangle - M_{\Phi_1}$ plane, induces substantial mixing between left-handed and right-handed sneutrinos. Sizeable mixing of this sort opens the intriguing possibility that the lightest sneutrino could be a legitimate dark matter candidate [22, 23, 24, 25] in Dirac leptogenesis, a possibility which would be interesting to investigate in the future.

The primary message of figure 2 is that the combined constraints from leptogenesis and flavor violation do not exclude Dirac leptogenesis in theories with low-scale sfermion masses. In general, the latter set of constraints tend to rule out theories with exceptionally high masses for the decaying particles Φ_1 and $\overline{\Phi}_1$ and exceptionally low values for $\langle\chi\rangle$.[9] Given current bounds on lepton flavor violation, there is still a sizeable region of parameter space within which the scenario succeeds, even when m_s is as low as 100 GeV. It is of interest, however, that MEG and the next generation of lepton flavor violation experiments will be able to probe the vast majority of this region, and the data from these experiments will be crucial in determining the viability of Dirac leptogenesis with weak-scale superpartners.

As was mentioned in section 2, some new symmetry must be posited in order to obtain late neutrino mass generation. In Dirac leptogenesis, neutrino masses are the result of the scalar component of the χ superfield acquiring a VEV which breaks this new symmetry, producing a Goldstone boson or pseudo-Goldstone boson, depending on the nature of the symmetry. Constraints on such bosons arise from both BBN and cosmic microwave background (CMB) considerations [26]

[9] Recall that the parameters $\langle\chi\rangle$, M_{Φ_1} and f enter in the effective neutrino Yukawa coupling given by 2, so when one of them is increased, we have to appropriately tune the other two to maintain the neutrino Yukawas at a fixed value. Such an adjustment, in turn, changes the different contributions from these parameters to off-diagonal entries in the slepton mass matrices.

as well as from the detection of abnormalities in the neutrino flux associated with supernova events [27], and they can become problematic (depending on the mass of the Goldstone boson) when the symmetry-breaking VEV is less than around 1 GeV. The value of $\langle\chi\rangle$ required by leptogenesis constraints is around 10 GeV, and thus the cosmological complications associated with breaking the additional symmetry necessary for Dirac leptogenesis do not pose any problem for the theory.

4 Conclusion

It has already been shown [3] that Dirac leptogenesis stands as a phenomenologically viable alternative to the standard leptogenesis picture. In theories with comparatively light scalars, however, flavor violation becomes a serious concern commonly solved by the Universality assumption. The Dirac leptogenesis superpotential (1) gives rise to new mass terms for sneutrinos and charged sleptons which induce mixings between flavor eigenstates after the breaking of both the electroweak symmetry and the new symmetry responsible for late neutrino mass generation. Experimental limits on flavor violation in heavy lepton decays such as $\mu \to e\gamma$ significantly constrain any theory which permits slepton flavor mixing, and in Dirac leptogenesis these constraints translate into bounds on the theory parameters M_{Φ_1} (the mass of the heavy decaying particle) and $\langle\chi\rangle$ (the VEV of the exotic scalar field). Baryogenesis requirements also place significant constraints on both of these parameters, and thus the question as to whether leptogenesis can be made to work at all when scalar masses are light is a highly nontrivial one.

In this work, we have shown that even when the masses of supersymmetric scalars are small, substantial regions of parameter space exist for which Dirac leptogenesis succeeds in producing a realistic baryon asymmetry for the universe while respecting current bounds on flavor violation. This is true even when the masses of supersymmetric particles are as low as ~ 100 GeV. Interestingly enough, for such light fields, which presumably can be discovered at the LHC, it is generically predicted that experiments such as MEG should have enough sensitivity to observe flavor changing effects if Dirac leptogenesis is in fact responsible for the baryon asymmetry of the universe.

5 Acknowledgments

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A Effective Couplings

For completeness, we list here the results used in our analysis for lepton flavor violating processes. The amplitudes A^L and A^R in equation (17) were computed in [13] and, with a trivial extension to include three right-handed sneutrinos, are given by

$$A^L = A^{(c),L} + A^{(n),L} \qquad \text{and} \qquad A^R = A^{(c),R} + A^{(n),R}, \tag{23}$$

where the individual amplitudes $A^{(c),L}$, $A^{(n),L}$, $A^{(c),R}$, and $A^{(n),R}$ are

$$\begin{aligned}
A^{(n)L} &= \frac{1}{32\pi^2}\sum_{A=1}^{4}\sum_{X=1}^{6}\frac{1}{m^2_{\tilde{\ell}_X}}\left[N^L_{iAX}N^{L*}_{jAX}\frac{1}{6(1-x_{AX})^4}\right.\\
&\quad\times(1-6x_{AX}+3x^2_{AX}+2x^3_{AX}-6x^2_{AX}\ln x_{AX})\\
&\quad\left.+N^L_{iAX}N^{R*}_{jAX}\frac{M_{\tilde{\chi}^0_A}}{m_{l_j}}\frac{1}{(1-x_{AX})^3}(1-x^2_{AX}+2x_{AX}\ln x_{AX})\right], &(24)\\
A^{(c)L} &= -\frac{1}{32\pi^2}\sum_{A=1}^{2}\sum_{X=1}^{6}\frac{1}{m^2_{\tilde{\nu}_X}}\left[C^L_{iAX}C^{L*}_{jAX}\frac{1}{6(1-x_{AX})^4}\right.\\
&\quad\times(2+3x_{AX}-6x^2_{AX}+x^3_{AX}+6x_{AX}\ln x_{AX})\\
&\quad\left.+C^L_{iAX}C^{R*}_{jAX}\frac{M_{\tilde{\chi}^-_A}}{m_{l_j}}\frac{1}{(1-x_{AX})^3}(-3+4x_{AX}-x^2_{AX}-2\ln x_{AX})\right], &(25)\\
A^{(n,c)R} &= A^{(n,c)L}|_{L\leftrightarrow R}. &(26)
\end{aligned}$$

Here, the indices A and X respectively label the gaugino (chargino or neutralino) and slepton (sneutrino or charged slepton) mass eigenstates, $x_{AX} \equiv m^2_{\chi A}/m^2_{\phi_X}$, and $C^{L,R}_{iAX}$ ($N^{L,R}_{iAX}$) denote the effective couplings of charged lepton i to chargino (neutralino) A and sneutrino (charged slepton) X. The flavor mixing terms in (**??**) enter into the overall rate (17) through $C^{L,R}_{iAX}$ and $N^{L,R}_{iAX}$, which contain elements of the matrices U_ν and U_ℓ that diagonalize the mass-squared matrices for sneutrinos and charged sleptons, respectively. The slepton masses also enter into the partial amplitudes (24-26).

The effective couplings $N^{L,R}_{iAX}$ and $C^{L,R}_{iAX}$ are

$$\begin{aligned}
N^R_{iAX} &= -\frac{g_2}{\sqrt{2}}\left([-(U_N)_{A2}-(U_N)_{A1}\tan\theta_W]U^\ell_{X,i}+\frac{m_{l_i}}{m_W\cos\beta}(U_N)_{A3}U^\ell_{X,i+3}\right),\\
N^L_{iAX} &= -\frac{g_2}{\sqrt{2}}\left(\frac{m_{l_i}}{m_W\cos\beta}(U_N)_{A3}U^\ell_{x,i}+2(U_N)_{A1}\tan\theta_W U^\ell_{X,i+3}\right),\\
C^R_{iAX} &= -g_2(O_R)_{A1}U^\nu_{X,i}, \quad \text{and}\\
C^L_{iAX} &= g_2\frac{m_{l_i}}{\sqrt{2}m_W\cos\beta}(O_L)_{A2}U^\nu_{X,i} &(27)
\end{aligned}$$

in terms of the chargino mixing matrices $(O_R)_{A,i}$ and $(O_L)_{A,i}$ the neutralino mixing matrix $U^N_{X,i}$, and the sneutrino and charged slepton mixing matrices $U^\nu_{X,i}$ and $U^\ell_{X,i}$. The chargino mixings matrices are defined by the relation

$$M_c^{diag} = (O_R) M_c (O_L)^T, \tag{28}$$

where

$$M_c = \begin{pmatrix} 0 & X \\ X^T & 0 \end{pmatrix}, \quad \text{where} \quad X = \begin{pmatrix} M_2 & \sqrt{2} M_W \cos\beta \\ \sqrt{2} M_W \sin\beta & \mu \end{pmatrix} \tag{29}$$

and M_c^{diag} is diagonal. The sneutrino mixing matrix $U^{\nu}_{X,i}$ and the charged slepton mixing matrix $U^{\ell}_{X,i}$ are defined by the relations

$$(m^2_{\tilde{\ell}^\pm})^{diag} = U^\ell m^2_{\tilde{\ell}^\pm} U^\dagger_\ell \quad , \quad (m^2_{\tilde{\nu}^\pm})^{diag} = U^\ell m^2_{\tilde{\nu}} U^\dagger_\ell, \tag{30}$$

where the matrices $m^2_{\tilde{\ell}^\pm}$ and $m^2_{\tilde{\nu}}$ are given by the sum of the MSSM contribution and the respective Dirac leptogenesis contributions in (8). The neutralino mixing matrix U_N is defined by the relation

$$(m_{\tilde{N}})^{diag} = U_N m_{\tilde{N}} U_N^\dagger, \tag{31}$$

where

$$m_{\tilde{N}} = \begin{pmatrix} M_1 & 0 & -M_Z \sin\theta_w \cos\beta & M_Z \sin\theta_w \sin\beta \\ 0 & M_2 & M_Z \cos\theta_w \cos\beta & -M_Z \cos\theta_w \cos\beta \\ -M_Z \sin\theta_w \cos\beta & M_Z \cos\theta_w \cos\beta & 0 & -\mu \\ M_Z \sin\theta_w \sin\beta & -M_Z \cos\theta_w \sin\beta & -\mu & 0 \end{pmatrix}. \tag{32}$$

B Boltzmann Equations Including Second Order Processes

In [3], we derived the system of Boltzmann equations for Dirac leptogenesis up to processes first order in the heavy, decaying particles Φ_1 and $\overline{\Phi}_1$ and showed that the dynamics of these fields could be expressed using only two equations: one for the lepton number L_{ϕ_Φ} in the heavy field sector and one for the abundance $Y^c_{\phi_\Phi}$ of the scalar component ϕ_1 of the Φ_1 supermultiplet. Here, we improve upon our previous calculations by including terms second order in ϕ_1. These terms only appear in the equation for the $Y^c_{\phi_\Phi}$ abundance, which becomes[10]

$$\frac{dY^c_{\phi_\Phi}}{dz} = -\gamma_D \left(\frac{Y^c_{\phi_\Phi}}{Y^{eq}_{\phi_\Phi}} - 1 \right) + \frac{1}{2} \gamma_L^{ID} \frac{L_{agg}}{Y^{eq}_{\phi_\Phi}} + \frac{1}{2} \gamma_R^{ID} \frac{L_{\nu_R}}{Y^{eq}_{\phi_\Phi}} - \gamma_A \left(\frac{(Y^c_{\phi_\Phi})^2}{(Y^{eq}_{\phi_\Phi})^2} - 1 \right), \tag{33}$$

where L_{ν_R} is the lepton number stored in sneutrinos, L_{agg} is the aggregate lepton number stored in the other, interacting fields that carry lepton number, and all abundances are normalized with

[10]Since the dynamics of scalar and fermion fields in Φ_1 and $\overline{\Phi}_1$ are assumed to be the same, contributions of the form $\tilde{\phi}_1 \phi_1 \to ij$ can be incorporated into γ_A.

respect to the entropy density s. The reaction densities γ_D, γ_L^{ID}, and γ_R^{ID} are given by

$$\gamma_D = \left(\frac{K_1(z)}{K_2(z)}\right) sY^c_{\phi_\Phi}\Gamma_D \tag{34}$$

$$\gamma_L^{ID} = \frac{1}{7}\frac{n^{eq}_{\phi_\Phi}}{n_\gamma}\left(\frac{K_1(z)}{K_2(z)}\right) sY^c_{\phi_\Phi}\Gamma_L \tag{35}$$

$$\gamma_R^{ID} = \frac{n^{eq}_{\phi_\Phi}}{n_\gamma}\left(\frac{K_1(z)}{K_2(z)}\right) sY^c_{\phi_\Phi}\Gamma R, \tag{36}$$

in terms of the Bessel functions $K_1(z)$ and $K_2(z)$, the photon number density n_γ, and the rates Γ_D, Γ_L, and Γ_R, which respectively represent the total decal width of the scalar component field ϕ_1 and the individual decay widths for the processes $\phi \to \tilde{\ell}\chi$ and $\phi \to \nu^c_R\tilde{H}^c_u$. The reaction density γ_A, which is associated with second order processes of the form $\phi_1\phi_1 \to ij$ and $\tilde{\phi}_1\phi_1 \to ij$, is given by

$$\gamma_A = \frac{T}{64\pi^4}\int_{s_{min}}^{\infty} s^{1/2}K_1\left(\frac{\sqrt{s}}{T}\right)\hat{\sigma}(s), \tag{37}$$

where T is temperature, s is the usual Mandelstam variable, and $\hat{\sigma}(s)$ is the total reduced cross section for annihilations of $\phi_1\phi_1$ and $\phi_1\tilde{\phi}_1$ into light fields. This is defined by the formula

$$\hat{\sigma}(s) = \frac{1}{8\pi s}\int_{t_-}^{t_+}\sum_i |\mathcal{M}_i(t)|^2 dt, \tag{38}$$

where both t and s denote the Mandelstam variables. The limits of integration are given by $t_\pm = M^2_{\Phi_1} - s(1 \mp r)/2$, with r defined below.

In supersymmetric Dirac leptogenesis, the total reduced cross-section γ_A, including all relevant decay processes, is

$$\begin{aligned}\hat{\sigma}^{tot}_{SUSY} &= \frac{1}{16\pi}\left[6g_Y^2g_2^2\left(\left(-7+\frac{4}{x}\right)r+\left(\frac{8}{x^2}-\frac{4}{x}+9\right)\ln\left(\frac{1+r}{1-r}\right)\right)\right.\\ &\quad +g_2^4\left(\left(32+\frac{66}{x}\right)r+3\left(-\frac{16}{x^2}-\frac{16}{x}+9\right)\ln\left(\frac{1+r}{1-r}\right)\right)\\ &\quad \left.+\,g_Y^4\left(\left(19-\frac{36}{x}\right)r+\left(\frac{16}{x^2}-\frac{8}{x}+17\right)\ln\left(\frac{1+r}{1-r}\right)\right)\right],\end{aligned} \tag{39}$$

where $x \equiv s/M_{\Phi_1}$, $r = \sqrt{1-4/x}$, and g_2 and g_Y are the $SU(2)$ and $U(1)_Y$ coupling constants. The calculation can also be performed for the non-supersymmetric case, where the result is

$$\begin{aligned}\hat{\sigma}^{tot}_{SM} &= \frac{1}{96\pi}\left[g_Y^2g_2^2\left(\left(36+\frac{144}{x}\right)r+144\left(\frac{2}{x^2}+\frac{1}{x}\right)\ln\left(\frac{1+r}{1-r}\right)\right)\right.\\ &\quad +3g_2^4\left(\left(39+\frac{196}{x}\right)r-144\left(\frac{2}{x^2}+\frac{3}{x}\right)\ln\left(\frac{1+r}{1-r}\right)\right)\\ &\quad \left.+\,g_Y^4\left(\left(53-\frac{116}{x}\right)r-\left(\frac{-96}{x^2}+\frac{48}{x}\right)\ln\left(\frac{1+r}{1-r}\right)\right)\right].\end{aligned} \tag{40}$$

The effect of such second order annihilation processes on the parameter space of Dirac leptogenesis is shown in figure 4, where, for comparison, we show two sets of leptogenesis exclusion contours:

one representing no second-order processes and one representing annihilation in a supersymmetric model. It is evident from this graph that second order processes do indeed lower the upper exclusion contour, though the effect is not a dramatic one.

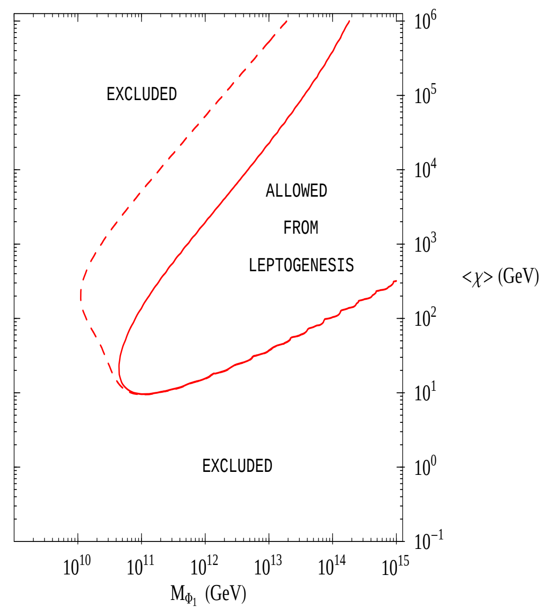

Figure 4: This figure illustrates the effect of second-order processes of the form $\phi_1\phi_1 \to ij$ and $\tilde{\phi}_1\phi_1 \to ij$ on the exclusion contours from leptogenesis. Contours are displayed for the case without annihilation (dashed line) and with annihilation (solid line).

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